

# Partial Breaking of Extended Supersymmetry

Jonathan A. Bagger <sup>\*</sup>

Department of Physics and Astronomy, Johns Hopkins University  
3400 N. Charles Street, Baltimore, MD 21218

In this talk we use nonlinear realizations to study the spontaneous partial breaking of rigid and local supersymmetry.

## 1. INTRODUCTION

In this talk we will consider the partial breaking of extended supersymmetry. For simplicity, we will restrict our attention to the case  $N = 2 \rightarrow N = 1$ , but much of what we find can be readily extended to the case of higher supersymmetries, spontaneously broken to  $N = 1$ .

There is a heuristic argument which implies that extended supersymmetry cannot be spontaneously broken to  $N = 1$  in four dimensions [1]. The argument runs as follows: Suppose that there are two supersymmetries, one broken and one unbroken. Since one supersymmetry is preserved, one supercharge must annihilate the vacuum. If the Hilbert space is positive definite, the supersymmetry algebra

$$\{Q_\alpha^A, \bar{Q}_{\dot{\alpha}B}\} = 2\sigma_{\alpha\dot{\alpha}}^m P_m \delta^A_B \quad (1)$$

implies that the Hamiltonian must also annihilate the vacuum. This, in turn, requires the second supercharge to annihilate the vacuum, so the other supersymmetry cannot be broken.

Hughes, Liu and Polchinski [2] found a legal loophole which allowed them to evade this argument. They exploited the fact that in theories with spontaneous symmetry breaking, the broken symmetry charges do not exist. This motivated them to consider the following current algebra,

$$\begin{aligned} \{Q_\alpha^1, \bar{J}_{m\dot{\alpha}1}\} &= 2\sigma_{\alpha\dot{\alpha}}^n T_{mn} \\ \{Q_\alpha^2, \bar{J}_{m\dot{\alpha}2}\} &= 2\sigma_{\alpha\dot{\alpha}}^n (v^4 \eta_{mn} + T_{mn}). \end{aligned} \quad (2)$$

Note that the right-hand sides of the two commutators differ by a constant. In the limit  $v \rightarrow 0$ , the

constant vanishes and the current algebra can be integrated to give the charge algebra (1). When  $v \neq 0$ , the current algebra cannot be integrated, and the second supersymmetry is spontaneously broken. (Of course, if there were just one supersymmetry, the constant would become the vacuum energy that signals spontaneous supersymmetry breaking.)

Hughes, Liu and Polchinski found an explicit realization of their algebra in terms of a four-dimensional supermembrane propagating in six-dimensional superspace. They found its invariant action and demonstrated that it realizes the partial breaking of extended supersymmetry.

The membrane approach leaves many open questions, some of which will be addressed in this talk. In particular, we would like to know if there are other realizations of partial supersymmetry breaking. We would like to know whether the  $N = 2$  supersymmetry gives rise to any restrictions on the  $N = 1$  matter couplings. And we would like to know if the system can be coupled to supergravity, because gravity can distinguish between the different stress-energy tensors on the right-hand side of eq. (2). (For alternative approaches to this subject, and other references, see [3] – [5], and references therein.)

## 2. COSET CONSTRUCTION

In this talk we will take a bottom-up approach to the subject of partial supersymmetry breaking. We will use nonlinear realizations to describe the effective  $N = 1$  theory which holds below the scale of the second supersymmetry breaking. We will use the formalism of Coleman, Wess and Zu-

<sup>\*</sup>Supported by the U.S. National Science Foundation, grant NSF-PHY-9404057.

mino [6], as modified by Volkov [7], to construct theories where the  $N = 1$  supersymmetry is manifest, and the second supersymmetry is nonlinearly realized.

The approach of Coleman, Wess, Zumino and Volkov is based on a coset decomposition of a symmetry group,  $G$ . We start with a group,  $G$ , of internal and spacetime symmetries, and partition the generators of  $G$  into three classes:

- $\Gamma_A$ , the generators of unbroken spacetime translations;
- $\Gamma_a$ , the generators of spontaneously broken internal and spacetime symmetries; and
- $\Gamma_i$ , the generators of unbroken spacetime rotations and unbroken internal symmetries.

The generators  $\Gamma_i$  close into the stability group,  $H$ .

Given  $G$  and  $H$ , we define the coset  $G/H$  in terms of an equivalence relation on the elements  $\Omega \in G$ ,  $\Omega \sim \Omega h$ , with  $h \in H$ . The coset can be thought of as a section of a fiber bundle with total space,  $G$ , and fiber,  $H$ .

This equivalence relation suggests that we parametrize the coset as follows,

$$\Omega = \exp iX^A \Gamma_A \exp i\xi^a(X) \Gamma_a. \quad (3)$$

Physically, the  $X^A$  play the role of generalized spacetime coordinates, while the  $\xi^a(X)$  are generalized Goldstone fields, defined on the generalized coordinates and valued in the set of broken generators  $\Gamma_a$ . There is one generalized coordinate for every unbroken spacetime translation, and one generalized Goldstone field for every spontaneously broken generator.

We define the action of the group  $G$  on the coset  $G/H$  by left multiplication,  $\Omega \rightarrow g\Omega = \Omega' h$ , with  $g \in G$ . In this expression,

$$\Omega' = \exp iX'^A \Gamma_A \exp i\xi'^a(X') \Gamma_a \quad (4)$$

and  $h = \exp i\alpha^i(g, X, \xi) \Gamma_i$ . The group multiplication induces nonlinear transformations on the coordinates  $X^A$  and the Goldstone fields  $\xi^a$ :

$$X^A \rightarrow X'^A, \quad \xi^a(X) \rightarrow \xi'^a(X'). \quad (5)$$

These transformations realize the full symmetry group,  $G$ . Note that the field  $\xi^a$  transforms by a

shift under the transformation generated by  $\Gamma_a$ . This confirms that  $\xi^a$  is indeed the Goldstone field corresponding to the broken generator  $\Gamma_a$ .

An arbitrary  $G$  transformation induces a compensating  $H$  transformation which is required to restore the section. This transformation can be used to lift any representation,  $R$ , of  $H$ , to a nonlinear realization of the full group,  $G$ , as follows,

$$\chi(X) \rightarrow \chi'(X') = D(h)\chi(X). \quad (6)$$

Here  $D(h) = \exp(i\alpha^i T_i)$ , where  $\alpha^i$  was defined below (4), and the  $T_i$  are generators of  $H$  in the representation  $R$ .

To construct an invariant action, it is helpful to have a vielbein, connection and covariant derivative, built from the Goldstone fields in the following way. One first computes the Maurer-Cartan form,  $\Omega^{-1}d\Omega$ , where  $d$  is the exterior derivative. One then expands  $\Omega^{-1}d\Omega$  in terms of the Lie algebra of  $G$ ,

$$\Omega^{-1}d\Omega = i(\omega^A \Gamma_A + \omega^a \Gamma_a + \omega^i \Gamma_i), \quad (7)$$

where  $\omega^A$ ,  $\omega^a$  and  $\omega^i$  are one-forms on the manifold parametrized by the coordinates  $X^A$ .

The Maurer-Cartan form transforms as follows under a rigid  $G$  transformation,

$$\Omega^{-1}d\Omega \rightarrow h(\Omega^{-1}d\Omega)h^{-1} - dh h^{-1}. \quad (8)$$

From this we see that the fields  $\omega^A$  and  $\omega^a$  transform covariantly under  $G$ , while  $\omega^i$  transforms by a shift. These transformations help us identify

$$\omega^A = dX^M E_M^A \quad (9)$$

as the covariant vielbein,

$$\omega^a = dX^M E_M^A \mathcal{D}_A \xi^a \quad (10)$$

as the covariant derivative of the Goldstone field  $\xi^a$ , and

$$\omega^i = dX^M \omega_M^i \quad (11)$$

as the connection associated with the stability group,  $H$ . With these building blocks, it is easy to construct an action invariant under the full group  $G$ .

The coset construction is very general and very powerful. For the case of internal symmetries, it

allows one to prove that any  $H$ -invariant action can be lifted to be  $G$ -invariant with the help of the Goldstone bosons. For  $N = 1$  supersymmetry, it can be used to show that any Lorentz-invariant action can be made supersymmetric with the help of the Goldstone fermion.

### 3. NONLINEAR SUPERSYMMETRY

In this section we will show that any  $N = 1$  supersymmetric theory can be made  $N = 2$  supersymmetric with the help of an  $N = 1$  Goldstone superfield. We will find that the Goldstone superfield can contain either an  $N = 1$  chiral or vector multiplet. Note, however, that the coset construction does not tell us anything about the underlying theory in which both supersymmetries are linearly realized. Indeed, such a theory is not even guaranteed to exist.

To begin, let us rewrite the  $N = 2$  supersymmetry algebra as follows,

$$\begin{aligned} \{Q_\alpha, \bar{Q}_{\dot{\alpha}}\} &= 2\sigma_{\alpha\dot{\alpha}}^a P_a, \quad \{S_\alpha, \bar{S}_{\dot{\alpha}}\} = 2\sigma_{\alpha\dot{\alpha}}^a P_a, \\ \{Q_\alpha, S_\beta\} &= 2\epsilon_{\alpha\beta} Z, \quad \{\bar{Q}_{\dot{\alpha}}, \bar{S}_{\dot{\beta}}\} = 2\epsilon_{\dot{\alpha}\dot{\beta}} \bar{Z}. \end{aligned} \quad (12)$$

Here  $Q_\alpha$  and  $S_\alpha$  are the supersymmetry generators,  $P_a$  the four-dimensional momentum operator, and  $Z$  is a complex central charge. In what follows, we will define  $Q_\alpha$  to be the unbroken  $N = 1$  supersymmetry generator, and  $S_\alpha$  to be its broken counterpart.

We shall first take a minimal approach, and choose the group  $G$  to be the supergroup whose algebra is (12). We will take the subgroup  $H$  to be the supergroup generated by  $P_a$ ,  $Q_\alpha$  and  $\bar{Q}_{\dot{\alpha}}$ . We parametrize the coset element  $\Omega$  as follows,

$$\begin{aligned} \Omega &= \exp i(x^a P_a + \theta^\alpha Q_\alpha + \bar{\theta}_{\dot{\alpha}} \bar{Q}^{\dot{\alpha}}) \\ &\quad \times \exp i(\psi^\alpha S_\alpha + \bar{\psi}_{\dot{\alpha}} \bar{S}^{\dot{\alpha}}). \end{aligned} \quad (13)$$

Here  $x$ ,  $\theta$  and  $\bar{\theta}$  are the coordinates of  $N = 1$  superspace, while  $\psi^\alpha$  and its conjugate  $\bar{\psi}_{\dot{\alpha}}$  are Goldstone  $N = 1$  superfields of (geometrical) dimension  $-1/2$ . These spinor superfields contain far too many component fields, so we need to find a set of consistent, covariant constraints to reduce the number of fields.

The correct constraints are most easily expressed in term of the  $N = 2$  covariant derivatives of the Goldstone superfield. These covariant

derivatives can be explicitly written as follows,

$$\begin{aligned} \mathcal{D}_\alpha &= D_\alpha - i(D_\alpha \psi \sigma^a \bar{\psi} + D_\alpha \bar{\psi} \sigma^a \psi) \omega_a^{-1m} \partial_m \\ \bar{\mathcal{D}}_{\dot{\alpha}} &= \bar{D}_{\dot{\alpha}} - i(\bar{D}_{\dot{\alpha}} \psi \sigma^a \bar{\psi} + \bar{D}_{\dot{\alpha}} \bar{\psi} \sigma^a \psi) \omega_a^{-1m} \partial_m \\ \mathcal{D}_a &= \omega_a^{-1m} \partial_m, \end{aligned} \quad (14)$$

where  $\omega_m^a \equiv \delta_m^a + i(\partial_m \psi \sigma^a \bar{\psi} + \partial_m \bar{\psi} \sigma^a \psi)$  and  $D_\alpha$ ,  $\bar{D}_{\dot{\alpha}}$  are ordinary flat  $N = 1$  superspace spinor derivatives. The covariant derivatives obey the following commutation relations,

$$\begin{aligned} \{\mathcal{D}_\alpha, \mathcal{D}_\beta\} &= -2i(\mathcal{D}_\alpha \psi^\gamma \mathcal{D}_\beta \bar{\psi}_{\dot{\gamma}} + (\alpha \leftrightarrow \beta)) \mathcal{D}_{\gamma\dot{\gamma}} \\ [\mathcal{D}_\alpha, \mathcal{D}_a] &= -2i(\mathcal{D}_\alpha \psi^\gamma \mathcal{D}_a \bar{\psi}_{\dot{\gamma}} + (\alpha \leftrightarrow a)) \mathcal{D}_{\gamma\dot{\gamma}} \\ \{\mathcal{D}_\alpha, \bar{\mathcal{D}}_{\dot{\beta}}\} &= 2i\sigma_{\alpha\dot{\beta}}^a \mathcal{D}_a - 2i(\mathcal{D}_\alpha \psi^\gamma \bar{\mathcal{D}}_{\dot{\beta}} \bar{\psi}_{\dot{\gamma}} \\ &\quad + (\alpha \leftrightarrow \dot{\beta})) \mathcal{D}_{\gamma\dot{\gamma}}, \end{aligned} \quad (15)$$

where  $\mathcal{D}_{\alpha\dot{\alpha}} \equiv \sigma_{\alpha\dot{\alpha}}^a \mathcal{D}_a$ .

The first set of constraints is simply [ 8]

$$\begin{aligned} \bar{\mathcal{D}} \mathcal{D} \psi_\alpha &= \mathcal{O}(\psi^3) \\ \mathcal{D}_\alpha \psi_\beta + \mathcal{D}_\beta \psi_\alpha &= \mathcal{O}(\psi^3). \end{aligned} \quad (16)$$

The right-hand side of this equation must be adjusted for consistency with (15). Remarkably, this can be done using the dimensionless invariants  $\bar{\mathcal{D}}_{\dot{\alpha}} \psi_\alpha$  and  $\mathcal{D}_\alpha \psi_\beta$  (together with their complex conjugates). It turns out that there is a unique, consistent solution order-by-order in powers of the Goldstone field.

The solution to the constraints (16) is easy to find in perturbation theory. To lowest order, it is just the chiral multiplet  $\phi$ ,

$$\begin{aligned} \psi_\alpha &= D_\alpha \phi + \mathcal{O}(\psi^3) \\ \bar{\mathcal{D}}_{\dot{\alpha}} \phi &= \mathcal{O}(\psi^3). \end{aligned} \quad (17)$$

In this expression,  $D_\alpha$  is the ordinary  $N = 1$  superspace spinor derivative.

The second set of constraints is [ 9]

$$\begin{aligned} \bar{\mathcal{D}}_{\dot{\alpha}} \psi_\alpha &= \mathcal{O}(\psi^3) \\ \mathcal{D}^\alpha \psi_\alpha + \bar{\mathcal{D}}_{\dot{\beta}} \bar{\psi}^{\dot{\beta}} &= \mathcal{O}(\psi^3). \end{aligned} \quad (18)$$

As above, the right-hand side must be adjusted for consistency with the algebra of covariant derivatives. Again, there is a unique, consistent solution. To lowest order in perturbation theory, it is just

$$\begin{aligned} \psi_\alpha &= W_\alpha + \mathcal{O}(\psi^3) \\ W_\alpha &= -\frac{1}{4} \bar{D} \bar{D} D_\alpha V + \mathcal{O}(\psi^3), \end{aligned} \quad (19)$$

where  $V$  is a real  $N = 1$  vector superfield. We see that the chiral and vector Goldstone multiplet can each be obtained to lowest order in perturbation theory. In fact, the consistency of the multiplets survives to all orders in perturbation theory.

The Goldstone action can be constructed order-by-order in the Goldstone fields. For the chiral case, it is simply [ 8]

$$S = v^4 \int d^4x d^2\theta d^2\bar{\theta} E [\phi^+ \phi + \mathcal{O}(\phi^4)]. \quad (20)$$

In this expression,  $E = \text{Ber}(E_M^A)$  is the superdeterminant of the vielbein, and  $v$  is the constant of dimension one which corresponds to the scale of the supersymmetry breaking. The action (20) is invariant under the full  $N = 2$  supersymmetry.

For the vector multiplet, the Goldstone action is just [ 9]

$$S = \frac{v^4}{4} \int d^4x d^2\theta \mathcal{E} W^2 + h.c. + \int d^4x d^4\theta E \mathcal{O}(W^4). \quad (21)$$

This action is invariant under  $N = 2$  supersymmetry. It is also gauge-invariant. Curiously enough, the gauge field contribution to the Goldstone action coincides with the expansion of the Born-Infeld action.

Having constructed the  $N = 2$  Goldstone action, we are now ready to add  $N = 2$  covariant matter. The basic ingredients are  $N = 2$  nonlinear generalizations of  $N = 1$  chiral and vector superfields. The generalized chiral superfields are defined by the constraint  $\bar{D}_{\dot{\alpha}}\chi = 0$ , while vector superfields are defined by the reality condition  $V = V^+$ . These constraints are consistent for either type of Goldstone multiplet.

The matter action is easy to write down for either Goldstone multiplet. The kinetic term is

$$S = \int d^4x d^4\theta E K(\chi^+, \chi) \quad (22)$$

while the superpotential term is

$$S = \int d^4x d^2\theta \mathcal{E} P(\chi). \quad (23)$$

As before,  $E$  and  $\mathcal{E}$  are superdeterminants of the supervielbein  $E_M^A$ . They can be adjusted to preserve the condition

$$\int d^4x d^4\theta E F(\chi) = 0. \quad (24)$$

This allows the matter action to be Kähler invariant, so the matter couplings are described in terms of Kähler manifolds, just as for  $N = 1$ .

It is not hard to generalize these results to include vector superfields. Our general conclusion is that any  $N = 1$  invariant theory can be lifted to be  $N = 2$  supersymmetric with the help of a Goldstone superfield. Furthermore, we find that the Goldstone superfield can contain either an  $N = 1$  chiral or vector multiplet.

Now that we have two explicit realizations of partial supersymmetry breaking, we can ask how they avoid the no-go argument discussed above. In each case, the nonlinear theory exploits the loophole of Hughes, Liu and Polchinski. The second supercurrent goes like  $J_\alpha^m \sim v^4 \sigma_{\alpha\dot{\alpha}}^m \bar{\lambda}^{\dot{\alpha}}$ , so its commutator with the second supercharge reproduces the algebra (2).

#### 4. GEOMETRY

The fact that the constraints need to be adjusted order-by-order in  $\psi_\alpha$  hints that a deeper structure underlies partial supersymmetry breaking. The  $N = 2$  supersymmetry does not provide enough symmetry to uniquely fix the covariant derivatives and the associated constraints. This intuition is borne out for the case of the chiral multiplet, where a much deeper set of symmetries acts on the Goldstone multiplet [ 8].

To see this, let us consider a coset where the group  $G$  contains not only  $N = 2$  supersymmetry, but also its maximal automorphism group,  $SO(5, 1) \times SU(2)$ , where the  $SU(2)$  acts on the two supersymmetry generators, and  $SO(5, 1)$  is the  $D = 6$  Lorentz group. (Under  $SO(5, 1)$ , the generators  $P_a$  and  $Z$  form a  $D = 6$  vector, while the supercharges form a  $D = 6$  Majorana-Weyl spinor). Let us take  $H$  to be  $SO(3, 1) \times SO(2) \times U(1)$ , where  $SO(3, 1) \times SO(2) \subset SO(5, 1)$ ,  $U(1) \subset SU(2)$ , and  $SO(3, 1)$  is the  $D = 4$  Lorentz group.

Our parametrization of the coset  $G/H$  involves the  $N = 1$  superspace coordinates, as well as dif-

ferent Goldstone superfields for each of the broken symmetries,

$$\begin{aligned}\Omega = & \exp i(x^a P_a + \theta^\alpha Q_\alpha + \bar{\theta}_{\dot{\alpha}} \bar{Q}^{\dot{\alpha}}) \\ & \times \exp i(\Phi Z + \bar{\Phi} \bar{Z} + \Psi^\alpha S_\alpha + \bar{\Psi}_{\dot{\alpha}} \bar{S}^{\dot{\alpha}}) \\ & \times \exp i(\Lambda^a K_a + \bar{\Lambda}^a \bar{K}_a + \Xi T + \bar{\Xi} \bar{T}).\end{aligned}\quad (25)$$

Here  $\Lambda^a, \bar{\Lambda}^a$  are the Goldstone superfields associated with the generators  $K_a, \bar{K}_a$  of  $SO(5,1)/SO(3,1) \times SO(2)$ . Similarly,  $\Xi, \bar{\Xi}$  are the Goldstone superfields for the broken generators  $T, \bar{T}$  of  $SU(2)/U(1)$ .

As before, the  $N = 1$  Goldstone superfields contain far more components than the minimal Goldstone multiplet. This motivates us to impose the following consistent set of constraints:

$$\begin{aligned}\bar{\mathcal{D}}_{\dot{\alpha}} \Phi = 0, \quad \mathcal{D}_\alpha \Phi = 0, \quad \mathcal{D}_a \Phi = 0 \\ \mathcal{D}_\alpha \Psi^\beta = 0, \quad \bar{\mathcal{D}}_{\dot{\alpha}} \Psi^{\dot{\beta}} = 0.\end{aligned}\quad (26)$$

These constraints allow us to express the Goldstone superfields  $\Psi^\alpha, \Lambda^a$  and  $\bar{\Xi}$  in terms of a single superfield  $\Phi$ . To lowest order, we find  $\Psi^\alpha = -\frac{i}{2} D^\alpha \Phi$ ,  $\Lambda_a = -\partial_a \Phi$ , and  $\bar{\Xi} = \frac{1}{4} D^2 \Phi$ . The constraint  $\bar{\mathcal{D}}_{\dot{\alpha}} \Phi = 0$  reduces  $\Phi$  to an  $N = 1$  chiral superfield.

The remarkable fact about this construction is that it reveals a geometrical role for each component of the chiral Goldstone multiplet. The scalar field,  $A$ , is the complex Goldstone boson associated with the spontaneously broken central charge symmetry. Its derivative,  $\partial_m A$ , is the Goldstone boson associated with  $SO(5,1)/SO(3,1) \times SO(2)$ . The  $F$ -component of  $\Phi$  is the complex Goldstone boson associated with the  $SU(2)/U(1)$ . Finally, the spinor is the Goldstone fermion that arises from the partially broken supersymmetry.

The action (20) turns out to be invariant under  $SO(5,1)$ , but it explicitly breaks  $SU(2)$  down to  $U(1)$ . Furthermore, any  $R$ -invariant  $N = 1$  matter action can be made  $SO(5,1)$  invariant. These facts hint that the Goldstone action might be related to the six-dimensional membrane of Hughes, Liu and Polchinski. Indeed, it is not hard to show that the chiral Goldstone action is precisely the gauge-fixed membrane action.

The geometry that underlies the vector case is presently under study. The Born-Infeld form of

the gauge action suggests that it might be related to some sort of D-brane in a higher dimension. But no matter what, one would like to find the Goldstone-type symmetries associated with the gauge field strength and the auxiliary field of the Goldstone multiplet.

In fact, the  $D$ -component of the Goldstone multiplet can be interpreted as the Goldstone boson associated with the following  $U(1)$  subgroup of the  $SU(2)$  automorphism symmetry:  $\delta\theta^\alpha = i\lambda\psi^\alpha$ ,  $\delta\psi^\alpha = i\lambda\theta^\alpha$ . Under such a transformation, the  $D$ -component is shifted by the constant parameter  $\lambda$ .

If we were to extend  $G$  in  $G/H$  by this  $U(1)$ , we would eliminate the dimensionless invariant  $\mathcal{D}^\alpha \psi_\alpha$  in favor of the corresponding Goldstone superfield. Even then, there would still be a dimensionless invariant associated with the gauge field strength,  $\mathcal{D}_{(\alpha} \psi_{\beta)}$ . This suggests that there is an extension of  $N = 2$  supersymmetry which associates a Goldstone-like symmetry with this field strength.

Moreover, gauge fields themselves can be interpreted as Goldstone fields associated with infinite-dimensional symmetry groups. This leads us to wonder whether the full symmetry of the new multiplet is some infinite-dimensional extension of  $N = 2$  supersymmetry.

## 5. SUPERGRAVITY

We have just seen that there are two possible Goldstone realizations of partial supersymmetry breaking. Both give rise to the current algebra (2). Because of the curious shift in the stress-energy tensor, one would like to couple the Goldstone multiplets to supergravity. Presumably this is possible [4], [5], in which case we would like to study the current algebra.

It is helpful to start our analysis by counting degrees of freedom. Since we wish to couple to  $N = 2$  supergravity, we need to include the  $N = 2$  supergravity multiplet, which contains one graviton, two gravitinos, and one vector. Moreover, since we are partially breaking supersymmetry, we also need to include enough fields to make up a massive  $N = 1$  spin-3/2 multiplet, which has one massive spin-3/2 field, two massive spin-1 fields,

and one massive spin-1/2 fermion. This counting suggests that we must include at least one chiral and one vector multiplet. We will only consider this minimal case in what follows.

The super-Higgs effect implies that after partial supersymmetry breaking, the massive gravitino eats the spin-1/2 Goldstone fermion. A priori, we do not know whether the Goldstone fermion is associated with the chiral or vector multiplet, so we will let it be a linear combination of the two fields.

From this starting point, it is possible to construct the most general supergravity coupling of the gravity and matter fields. Recently, we carried out this construction to lowest nontrivial order in  $1/v^2$  and  $\kappa$ , the inverse Planck mass [11]. We assumed that supersymmetry is partially broken, that the cosmological constant is zero, and that the second gravitino acquires a mass proportional to  $\kappa v^2$ .

In our calculation, we demanded that the action be invariant under  $N = 2$  and central charge symmetry, and we required the transformations to close (up to equations of motion on the fermions). In the end, we found the coupling to be unique. As a check, we verified that in unitary gauge, our action and transformations reduce to those of a massive  $N = 1$  spin-3/2 field, first worked out in ref. [10].

Our results show that the supergravity coupling of the Goldstone system fixes the Goldstone fermion to be a particular linear combination of the fermions from the chiral and vector multiplets. The second supercurrent receives contributions from both fermions, as well as from the second gravitino.

Once we derived the supercurrents, we used them to compute the current algebra. We found that the contribution from the second gravitino exactly cancels the contributions from the other two fermions, so the current algebra takes the following form,

$$\{Q_\alpha^A, \bar{J}_{m\dot{\alpha}B}\} = 2\sigma_{\alpha\dot{\alpha}}^n T_{mn} \delta^A_B. \quad (27)$$

In the presence of supergravity, there is no confusion about the stress-energy tensor. There is just one such tensor, and it shows up on the right-hand side of the current algebra.

How, then, is the no-go argument avoided? Our results indicate that the negative-norm components of the gravitino invalidate one of the main assumptions behind the argument. There is no obstacle to partial supersymmetry breaking in the presence of gravity. The connection between this result and the membranes/D-branes of rigid supersymmetry breaking is, at present, an urgent and open question.

It is a pleasure to thank my collaborators, Sasha Galperin and Sam Osofsky, for their insights on the work presented here. I would also like to express my debt to Victor Ogievetsky for teaching me the nonlinear approach to spontaneously broken spacetime symmetries.

## REFERENCES

1. E. Witten, Nucl. Phys. B188 (1981) 513.
2. J. Hughes, J. Liu and J. Polchinski, Phys. Lett. 180B (1986) 370; J. Hughes and J. Polchinski, Nucl. Phys. B278 (1986) 147.
3. I. Antoniadis, H. Partouche and T. Taylor, Phys. Lett. B372 (1996) 83; S. Ferrara, L. Girardello and M. Porrati, Phys. Lett. B376 (1996) 275.
4. M. Awada, M. Duff and C. Pope, Phys. Rev. Lett. 50 (1983) 294; M. Duff, B. Nilsson and C. Pope, Phys. Rep. 130 (1986) 1.
5. S. Ferrara, L. Girardello and M. Porrati, Phys. Lett. B366 (1996) 155; P. Fré, L. Girardello, I. Pesando and M. Trigiante, hep-th/9607032; M. Porrati, hep-th/9609073, and references therein.
6. S. Coleman, J. Wess and B. Zumino, Phys. Rev. 177 (1969) 2239.
7. D.V. Volkov, Sov. J. Particles and Nuclei 4 (1973) 3; V.I. Ogievetsky, Proceedings of X-th Winter School of Theoretical Physics in Karpacz, vol. 1, p. 227 (Wroclaw, 1974).
8. J. Bagger and A. Galperin, Phys. Lett. B336 (1994) 25.
9. J. Bagger and A. Galperin, hep-th/9608177, to appear in Phys. Rev. D.
10. S. Ferrara and P. van Nieuwenhuizen, Phys. Lett. B127 (1983) 70.
11. J. Bagger, A. Galperin and S. Osofsky, to appear.